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Abstract. The exact solution of the classical nonlinear equation of motion for a relativistic electron in the field of two electromagnetic (EM) waves is obtained. For the particular case of the linearly polarized standing EM wave in the planar optical cavity, the intensity of the nonlinear Compton scattering, the time of flight, and the momentum variation after the relativistic electron passes along the cavity axis are calculated in weak and strong field limits. The extent of these effects depends on the initial phase of the EM wave when the electron enters the cavity. This can be used for the production, diagnosis, and acceleration of relativistic electron (positron) microbunches.

INTRODUCTION

The theory of the Compton effect in interfering EM waves, in particular, in two counter-propagating plane waves, has been addressed previously to describe the Kapitza-Dirac effect [1,2], Compton lasers [3,4], and inverse Compton laser acceleration [2,5]. The physical principle of the nanometer-resolution Shintake electron beam profile monitor [6-9] is also based upon the understanding of the Compton effect in a standing EM wave. The vacuum beat wave laser accelerator concept [10] relies on the ponderomotive acceleration resulting from the beat wave produced by the interaction of two copropagating laser beams

Temporal diagnostics of ultra-fine electron microbunches (sized to a portion of the laser wavelength) is another potential application for intense standing EM waves. Production and reliable characterization of such microbunches are essential for development of far-field and near-field laser accelerator schemes into practically meaningful monochromatic electron (positron) accelerators. The example of such a scheme is the staged electron laser acceleration experiment, STELLA, at the Brookhaven Accelerator Test Facility [11]. In this experiment, a train of the 1 μm (3 fs) thin electron microbunches, produced by the IFEL method and periodical to the CO_2 laser wavelength, λ=10 μm, is phased to the inverse Cherenkov laser acceleration stage driven by the same CO_2 laser. Observation of Compton scattered radiation from the electron interacting with a standing EM wave, produced by two counter-propagating CO_2 laser beams, may permit a direct assessment of the microbunch temporal structure. The inverse process may provide also an alternative mechanism to generate microbunches starting with a quasi-continuous electron pulse.

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Theoretical studies of the Compton effect in two interfering EM waves, comprehensively reviewed by M.V. Fedorov [2], generally capitalize on various aspects of the perturbation theory or other approximate approaches. As long as we are interested in processes nonlinear to the field, the approach based on exact solution of equations of motion is the most appropriate. The exact solutions of classical equations of the electron motion in a single plane EM wave have been obtained by many authors using different methods (see [12-16] and references therein). A natural extension of this approach is to apply the same methods to the case of two interfering plane waves.

In the recently submitted publication [17], this intent is partially accomplished using the approach developed in [13, 14]. In the present paper, we review results from [17] which are relevant to the problem of the microbunch diagnostics. In Section 2, we show the exact solution of classical nonlinear equations of electron motion in the field of two plane EM waves.

The particular case of a standing wave is considered in Section 3. We calculate the time of flight for the electron passing through the radiation-filled plane-parallel optical resonator and demonstrate that this parameter depends upon the phase of the standing EM wave at the moment when the electron enters the cavity.

In Section 4, we calculate the phase dependent intensity radiated due to Compton scattering when the relativistic electron passes along the axis of the plane-parallel optical cavity. The obtained results, discussed in Section 5, can be used for short relativistic electron bunch diagnostics.

SOLUTION OF NONLINEAR EQUATION OF MOTION

Following the approach outlined in [13,14], we consider a classical equation of the electron motion in EM field:

$$\frac{d\pi_\mu}{d\tau}=\frac{e}{m}F_{\mu\nu}\pi_\nu, \tag{1}$$

where π_μ is the four-dimensional vector of the electron energy-momentum $\pi_\mu(\varepsilon,\vec{p})$ equal to $\pi_\mu^2=\varepsilon^2-p^2$, τ is time in the relativistic frame of the electron (local or proper time), and $F_{\mu\nu}$ is a tensor of EM field, expressed through the vector potentials $A_{\nu,\mu}$,

$$F_{\mu\nu}=\frac{\partial A_\nu}{\partial x_\mu}-\frac{\partial A_\mu}{\partial x_\nu}, \tag{2}$$

$$A_\nu=a_\nu^{(1)}(\varsigma_1)+a_\nu^{(2)}(\varsigma_2), \tag{3}$$

where $a_\nu^{(1)}(\varsigma_1)$, $a_\nu^{(2)}(\varsigma_2)$ are vector potentials of two linearly polarized plane waves with frequencies ω_1 and ω_2 , $\varsigma_1=k_\mu^{(1)}x_\mu$ and $\varsigma_2=k_\mu^{(2)}x_\mu$ are the phases of corresponding plane waves, and $k_{\mu}^{(1)}(\omega_1, \vec{\omega}_1)$ and $k_{\mu}^{(2)}(\omega_2, \vec{\omega}_2)$ are four-dimensional wave vectors. In Eqs. (1), (2) and below, the system of units with $c = 1, h = 1$ is used.

A four-dimensional electron vector x_{μ} is

$$x_{\mu}(\tau) = \frac{1}{m}\int_{\tau_i}^{\tau} \pi_{\mu}(\tau)d\tau + x_{\mu}(\tau_i). \tag{4}$$

Consider linearly polarized waves that satisfy generalized transverse conditions

$$\left(k^{(1)}a^{(2)}\right) = 0, \quad \left(k^{(2)}a^{(1)}\right) = 0 \ . \tag{5}$$

We search for the solution of Eqs. (1)-(4) as a linear decomposition over four-vectors $p_{\mu}, a_{\mu}^{(1)}, a_{\mu}^{(2)}, k_{\mu}^{(1)}, k_{\mu}^{(2)}$,

$$\pi_{\mu}(\tau) = p_{\mu} - e\left[a_{\mu}^{(1)}(\varsigma_1) + a_{\mu}^{(2)}(\varsigma_2)\right] + k_{\mu}^{(1)} f_1(\tau) + k_{\mu}^{(2)} f_2(\tau), \tag{6}$$

where p_{μ} is the four-vector of the electron initial momentum before the electron enters the EM field. This brings us to the following result [17]:

$$f_{1,2} = \frac{\left(k^{(2,1)}p\right)}{k^{(1)}k^{(2)}}\left[\exp\left(\int_{\tau_i}^{\tau} F_{1,2}d\tau\right) - 1\right], \tag{7}$$

where quantities F_1 and F_2 are

$$F_{1,2} \equiv \frac{e\dfrac{d}{d\tau}\left[\left(a^{(1,2)}p\right) - \dfrac{e}{2}\left(a^{(1,2)}\right)^2\right] - e^2 a^{(2,1)}\dfrac{da^{(1,2)}}{d\tau}}{e\left(\left(a^{(1)} + a^{(2)}\right)p\right) - \dfrac{e^2}{2}\left(a^{(1)} + a^{(2)}\right)^2 + \dfrac{\left(k^{(1)}p\right)\left(k^{(2)}p\right)}{k^{(1)}k^{(2)}}} \ , \tag{8}$$

and the phases of the fields satisfy the equation

$$d\varsigma_{1,2} = \frac{\left(k^{(1,2)}p\right)}{m}\exp\left(\int_{\tau_i}^{\tau} F_{2,1}d\tau\right) \ . \tag{9}$$

ELECTRON IN PLANE-PARALLEL OPTICAL RESONATOR

Let us apply the solution obtained in Section 2 to the particular case of a standing EM wave formed by two linearly polarized plane waves of the same frequency ω, counter-propagating along the *x*-axis and polarized along the *y*-axis with

$$
\begin{aligned}
&a_y^{(1)}(\varsigma_1) = -(E_0/2\omega)\cos\varsigma_1 \ , \ \varsigma_1 = \omega(t-x) \ , \\
&a_y^{(2)}(\varsigma_2) = (E_0/2\omega)\cos\varsigma_2 \ , \ \varsigma_2 = \omega(t+x) \ , \\
&A_y = a_y^{(1)} + a_y^{(2)} = -(E_0/\omega)\sin\omega t \sin\omega x \ ,
\end{aligned}
\tag{10}
$$

where E_0 is the amplitude of electric field. A standing wave is confined between two conducting surfaces (mirrors) placed at x=0 and x=L=$n\lambda/2 = n\pi/\omega$, n=1,2,3....

The electron moving along the axis of the optical resonator enters the cavity at the moment $t = t_i(\tau_i)$ and leaves it at $t = t_f(\tau_f)$ (t is time in the laboratory scale, τ is the local time of the moving electron). Time of flight in the lab system is

$$
T = t_f - t_i = \int_{\tau_i}^{\tau_f} \gamma d\tau = \frac{1}{m}\int_{\tau_i}^{\tau_f} \pi_0 d\tau , \tag{11}
$$

where γ is a Lorentz factor of the electron.

Weak Fields

Let us consider the case of the weak field, when the normalized field amplitude is $\eta^2 = e^2a^2/m^2 \langle\langle 1$, where $a = E_0/2\omega$ is the amplitude of the potentials in Eq. (10). In this case, Eq. (8) for $F_{1,2}$ can be expanded in series of η^2, and Eq. (7) for $f_{1,2}$ takes the form

$$
\begin{aligned}
f_{1,2} \approx & \frac{e}{\left(k^{(1,2)}p\right)}\left[\left(a^{(1,2)}p\right) - \frac{e}{2}\left(a^{(1,2)}\right)^2\right] - \\
& \frac{e}{\left(k^{(1,2)}p\right)}\left\{\left[(a^{(1,2)}p) - \frac{e}{2}(a^{(1,2)})^2\right]_i - e\int_{\tau_i}^{\tau} a^{(2,1)} \frac{da^{(1,2)}}{d\tau} d\tau\right\}.
\end{aligned}
\tag{12}
$$

In the considered below ultrarelativistic case when $\vec{p}$ is collinear to $\vec{k}^{(1)}$, we obtain: $k^{(1)}p = \omega(\varepsilon - p) \approx \omega\,\varepsilon m^2/2p^2$, $\quad k^{(2)}p = \omega(\varepsilon + p) \approx 2\omega\,\varepsilon,\quad$ and $|f_2|\langle\langle|f_1|$.

By Eq. (12),

$$
f_{1f} = -\frac{2p^2\eta^2}{\omega}\left\{\frac{1}{2}\left(\cos^2\varsigma_{1f} - \cos^2\varsigma_{1i}\right) - \int_{\varsigma_{1i}}^{\varsigma_{1f}} \cos\varsigma_2 \sin\varsigma_1 d\varsigma_1\right\}, \tag{13}
$$

where

$$
\varsigma_{1i} = \omega t_i \quad (t = t_i\,,\ x = 0), \quad \varsigma_{1f} = \omega(t_i + T) - \omega L \ , \tag{14}
$$

L is the length of the resonator, and T is the time of flight,

$$T \equiv t_f - t_i = \frac{1}{m}\int_{\tau_i}^{\tau_f} \pi_0 d\tau \approx \frac{p}{m}\left(\tau_f - \tau_i\right) + \frac{\omega}{m}\int_{\tau_i}^{\tau_f} f_1 d\tau ,$$

$$L = \frac{p}{m}\left[1 - \frac{p\eta^2}{\varepsilon}\left(\frac{1}{2} - \cos^2\varsigma_{1i}\right)\right]\left(\tau_f - \tau_i\right) . \tag{15}$$

Neglecting in Eq. (13) terms proportional to $\left(k^{(1)}p\right)/\left(k^{(2)} \pm k^{(1)}, p\right) \sim 1/4\gamma^2$, we arrive to

$$f_{1f} = -\frac{p^2\eta^2}{\omega}\left(\cos^2\varsigma_{1f} - \cos^2\varsigma_{1i}\right) . \tag{16}$$

Expressing T and L through the local time interval $\left(\tau_f - \tau_i\right)$ we obtain:

For a short resonator, when $k^{(1)}p\left(\tau_f - \tau_i\right)/m \approx \omega L/2\gamma^2 \ll 1, \qquad (\omega L = n\pi)$,

$$T = \varepsilon\left(\tau_f - \tau_i\right)/m, \quad L = p\left(\tau_f - \tau_i\right)/m, \quad T = \varepsilon L/p . \tag{17}$$

For a long resonator, when $k^{(1)}p\left(\tau_f - \tau_i\right)/m \approx \omega L/2\gamma^2 \gg 1$,

$$T \approx \frac{\varepsilon}{m}\left[1 - \frac{p^2\eta^2}{\varepsilon^2}\left(\frac{1}{2} - \cos^2\varsigma_{1i}\right)\right]\left(\tau_f - \tau_i\right) ,$$

$$L = \frac{p}{m}\left[1 - \frac{p\eta^2}{\varepsilon}\left(\frac{1}{2} - \cos^2\varsigma_{1i}\right)\right]\left(\tau_f - \tau_i\right) , \tag{18}$$

or

$$T = \frac{\varepsilon}{p}L\left[1 + \frac{m^2\eta^2}{2\varepsilon p}\left(\frac{1}{2} - \cos^2\varsigma_{1i}\right)\right] . \tag{19}$$

In any case, the final phase, ς_{1f}, is

$$\varsigma_{1f} \approx \varsigma_{1i} + \omega L(\varepsilon - p)/p \approx \varsigma_{1i} + \omega L/2\gamma^2 . \tag{20}$$

We can define the coherent interaction distance, d_c, which corresponds to the distance where the electron phase slippage relative to the copropagating wave is $\Delta\varsigma_1 = \varsigma_1 - \varsigma_{1i} = \pi$ (compare with coherent radiation distance or radiation formation zone [18]). Then, according to Eq. (20), $d_c = \frac{2\pi}{\omega}\gamma^2 = \lambda\gamma^2$ and definitions for "short" and "long" resonators can be modified to $L \ll d_c$ for the short resonator, and $L \gg d_c$ for the long resonator.

Strong Fields

When trying to consider the strong field case, $\eta^2 = e^2a^2/m^2 >> 1$, we realize that straightforward calculation of the integrals in Eqs. (7) and (8) is difficult due to rapid variations of the expressions under the integrals. By transformation and integration of general expressions for $f_{1,2}$ the exact equation valid for arbitrary orientation of $\vec{p}$ and $\vec{k}^{1,2}$ and for an arbitrary field strength, η, can be obtained [17]. For counter-propagating plane waves, $k_\mu^{(1)}$ and $k_\mu^{(2)}$, linearly polarized along the y-axis, and with $\vec{p}$ directed along $k_1(\vec{k}_x,0,0)$, the exact equation takes the form

$$\frac{2\omega}{\varepsilon+p}f_1+\frac{2\omega}{\varepsilon-p}f_2+\frac{4\omega^2}{m^2}f_1f_2=\eta^2X^2, \tag{21}$$

where the condition $\varsigma_{1i}=\varsigma_{2i}$ is taken into account and

$$X^2 \equiv (\cos\varsigma_1-\cos\varsigma_2)^2. \tag{22}$$

By applying Eq. (21) to the case of $\eta^2 << 1$, that corresponds to $E_0 \to 0$ (or $\omega \to \infty$), we obtain

$$f_{1,2}=m^2\eta^2y_{1,2}/\left(k^{(1,2)}p\right), \tag{23}$$

where $y_{1,2}$ are functions of $\varsigma_{1,2}$ and $y_1+y_2=X^2$.

For the strong field case, $\eta^2 >> 1$, we use similar arguments as for $\eta^2 << 1$ above. First, notice that the condition $\eta^2 >> 1$ is fulfilled when $E_0 \to \infty$ (or $\omega \to 0$). Applying the limit $\eta^2 \to \infty$ to left and right sides of Eq. (21), we find the solution

$$f_{1,2}=m^2\eta y_{1,2}/2\left(k^{(1,2)}p\right) \;, \tag{24}$$

where $y_{1,2}$ are functions of $\varsigma_{1,2}$ and $y_1y_2=X^2$. The simplest choice that provides the symmetry condition, $k^{(1)} \leftrightarrow k^{(2)}$, $f_1 \leftrightarrow f_2$, is $y_1=y_2=|X|$ as is adopted below.

COMPTON SCATTERING IN STANDING EM WAVE

The total radiated energy during the passage of the electron through the optical resonator of the length L is

$$\Delta E = \frac{2}{3} e^2 \int_{t_i}^{t_f} w^2 dt = \frac{2}{3} e^2 m \int_{\tau_i}^{\tau_f} w^2 \pi_0 d\tau , \tag{25}$$

where w is the four- vector of acceleration,

$$w^2 = \frac{1}{m^2} \frac{d\pi_\mu}{d\tau} \frac{d\pi_\mu}{d\tau} . \tag{26}$$

For the weak field case, $\eta^2 << 1$, π_μ is given by Eqs. (6)-(8), and (23). For the strong field case, use Eq. (24) instead of Eqs. (23).

Consider the weak field approximation for the ultrarelativistic electron motion along the standing wave axis. Then, the expression under the integral in Eq. (25), up to the terms proportional to η^4, is

$$w^2 \pi_0 \approx \frac{e^2 \varepsilon}{m^4} \left(\frac{da^{(2)}}{d\varsigma_2} \right)^2 \left\{ \left(k^{(2)} p \right) + f_1 \left[2 \left(k^{(2)} p \right) \left(k^{(1)} k^{(2)} \right) + \frac{\omega}{\varepsilon} \left(k^{(2)} p \right) \right] \right\} . \tag{27}$$

The integral in Eq. (25) looks slightly different for long and short cavities:

$$\Delta E = \frac{1}{3} \frac{e^2 \eta^2 \omega^2 \varepsilon (\varepsilon + p)^2}{m^3} \left(\tau_f - \tau_i \right) \left\{ 1 + 3\eta^2 \varphi(\varsigma_{1i}) \right\} , \tag{28}$$

$$\varphi(\varsigma_{1i}) = \begin{cases} 1/2 - \cos^2 \varsigma_{1i} , & L\omega/2\gamma^2 \rangle\rangle 1 ; \\ 1/2 - \cos^2 \varsigma_{1i} + \cos\varsigma_{1i} , & L\omega/2\gamma^2 \langle\langle 1 . \end{cases} \tag{29}$$

Using Eqs. (17)-(19) for the time of flight, T, we obtain the following expressions for the average intensity, $I = \Delta E/T$, of the scattered radiation:

for a long resonator

$$I_{long} \approx \left[e^2 \eta^2 \omega^2 (\varepsilon + p)^2 / 3m^2 \right] \left[1 + 4\eta^2 \left(1/2 - \cos^2 \varsigma_{1i} \right) \right] , \tag{30}$$

and for a short resonator

$$I_{short} \approx \left[e^2 \eta^2 \omega^2 (\varepsilon + p)^2 / 3m^2 \right] \left[1 + 3\eta^2 \left(1/2 - \cos^2 \varsigma_{1i} + \frac{1}{2} \cos\varsigma_{1i} \right) \right] . \tag{31}$$

In the CGS system, the right hand side of Eqs. (28) and (29) need to be multiplied by the factor c^{-7}, the right hand side of Eqs. (30) and (31) - multiplied by the factor c^{-5}, and p will be replaced by pc.

Eqs. (30) and (31) comply, to the precision of up to η^2, with the results obtained for a single plane wave in Ref. [13], [14] (if we take $\vec{p} = 0$). The difference is due to the terms proportional to η^4 which depend upon the initial phase. The dependence of the nonlinear Compton scattering on phase can be used for the electron microbunch diagnostics.

For the ultrarelativistic strong field case, when $\eta^2 >> 1$, the total radiated energy during the electron passage through the optical resonator is

$$\Delta E \approx \left(16\varepsilon^4\eta^4 e^2\omega/3m^4\right)\int_{\varsigma_{1i}}^{\varsigma_{1f}} \sin^2\varsigma_2\left(\cos\varsigma_1 - \cos\varsigma_2\right)d\varsigma_1 \approx \left(2\varepsilon^4\eta^4 e^2\omega/3m^4\right)\left[4\left(\varsigma_{1f} - \varsigma_{1i}\right) + \sin 2\varsigma_{1f} - \sin 2\varsigma_{1i}\right]. \quad (32)$$

In Eq. (32), $\sin^2\varsigma_2$ and $\cos^2\varsigma_2$ are set equal to their average value 1/2, that is possible due to rapid variation of ς_2.

For a short resonator ($\omega L/2\gamma^2 << 1$), Eq. (32) can be reduced to

$$\Delta E = 2\varepsilon^2\eta^4 e^2\omega^2 L\left[2 + \cos 2\varsigma_{1i}\right]/3m^2; \quad (33)$$

and for a long resonator ($\omega L/2\gamma^2 >> 1$)

$$\Delta E \approx 8e^2\varepsilon^2\eta^4\omega^2 L/3m^2. \quad (34)$$

In the CGS system, the right hand side of Eqs. (33) and (34) is multiplied by c^{-6}.

DISCUSSION AND CONCLUSIONS

Following the classical approach developed in [13,14] for a single planar EM wave, the general solutions for electron motion in the field of two linearly polarized planar EM waves are obtained. Special consideration is given to the problem of radiation of the ultrarelativistic electrons, propagating along the standing EM wave axis.

The total radiated energy due to Compton scattering is defined primarily by the wave counter-propagating to the direction of the electron momentum. This feature is physically understandable, if we take into account that the relativistic electron experiences quickly oscillating force from the counter-propagating component of the standing wave that causes the electron to radiate, but moves practically in phase with the co-propagating component. Constant amplitudes of the electric and magnetic fields produced by the co-propagating wave define a drift of the electron trajectory in the cavity. As a result, the amplitude of the electron oscillation is modulated by the initial phase, and this is revealed through radiation.

We can expect that the overall angular and spectral spread of the radiated photons shall generally obey distributions obtained in [13,14,19,20] for backscattering of a single laser beam on the relativistic electron beam. However more detailed analysis may be necessary, especially for the spectral distribution of the phase dependent terms in Eqs. (19, 28-29, 33).

Phase dependence of the nonlinear Compton scattering in a standing laser wave may be used for microbunch characterization (bunch duration, longitudinal charge distribution). This is required for the advanced laser acceleration experiments, such as the STELLA experiment at the Brookhaven ATF [11].

Propagating along the axis of the standing EM wave, the electron can loose its energy via Compton scattering or acquire it in the inverse process.

Using general solutions obtained in Section 3 for the electron passing through the radiation filled plane-parallel cavity, we can address a problem of electron acceleration in vacuum. For weak field case, $\eta \ll 1$, the momentum of the electron at the exit from the cavity is

$$p_f = p_i - \frac{p^2\eta^2}{\varepsilon}\left(\cos^2\varsigma_{1f} - \cos^2\varsigma_{1i}\right) = p_i\left\{1 - \frac{p\eta^2}{\varepsilon}\left[\cos^2\left(\varsigma_{1i} + \frac{L\omega}{2\gamma^2}\right) - \cos^2\varsigma_{1i}\right]\right\}. \quad (35)$$

For a short resonator ($\omega L/2\gamma^2 \ll 1$)

$$p_f = p_i\left[1 - \frac{p\eta^2}{\varepsilon}\left(\frac{1}{2} - \cos^2\varsigma_{1i} - \sin 2\varsigma_{1i}\right)\right], \quad (36)$$

and for a long resonator ($\omega L/2\gamma^2 \gg 1$), after averaging over small changes in frequency, $\frac{\Delta\omega}{\omega} \langle\langle \frac{4\gamma^2}{\omega L} \langle\langle 1$,

$$p_f = p_i\left[1 - \frac{p\eta^2}{\varepsilon}\left(\frac{1}{2} - \cos^2\varsigma_{1i}\right)\right]. \quad (37)$$

Thus, if $\cos^2\varsigma_{1i} \rangle 1/2$, the acceleration takes place. For example, if $\varsigma_{1i} = \omega t_i = 0$ and η^2=0.25, then, after passing the resonator, p_f=1.12 p_i.

Contrary to Compton scattering, electron acceleration is primarily due to the component of the standing wave collinear with the electron propagation. The effect of the second counter-propagating wave is dumped due to the phase averaging. In the classical approach considered here, the acceleration of electrons is due to the ponderomotive force related to the linear and nonlinear parts of the Lorentz force as it is typical for the ponderomotive acceleration processes [10, 15, 16].

The facts that the electron-laser interaction is localized within the finite optical cavity and the accelerating force is nonlinear to the laser field circumvent the Lawson-Woodward theorem [21], that otherwise forbids the residual electron energy gain from EM waves in vacuum.

The revealed dependence of the electron momentum at the exit of the optical resonator upon the initial phase of the standing EM wave may be potentially used for electron (positron) microbunch diagnostics as well.

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